

Pressure Corrections to the Equation of State in the Nuclear Mean Field.

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We show the connection between stiffness of equation of state in a Relativistic Mean Field (RMF) of Nuclear Matter (NM) and the existence of a strong violation of longitudinal Momentum Sum Rule (MSR) in RMF for a finite pressure. The increasing pressure between nucleons starts to increase the ratio of nucleon Fermi to average single particle energy and according to the Hugenholtz-van Hove theorem valid for NM the MSR is broken. In order to satisfy that MSR we propose changes which modify the nucleon Parton Distribution Function (PDF) above a saturation density. The course of Equation of State in our modified RMF model is very close to semi-empirical estimation and to results obtained from extensive DBHF calculations with a Bonn A potential. Other features of the model includes a good values saturation properties including spin-orbit term. The proper stiffness of EoS recently discussed in an application to compact and neutron stars is important when studying star properties (mass-radius constraint), especially the mass of "PSR J16142230" the most massive known neutron star, which rules out many soft equations of state including exotic matter. An admixture of additional hyperons are discussed in our approach.

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I. INTRODUCTION

Experimentally, in Deeply Inelastic electron Scattering (DIS) on nuclear targets, fotons with large negative momentum square $-q^2 = Q^2 > 1\text{GeV}^2$ and large energy transfer ν , probe interacting hadrons - a kind of moving sub-targets. Start with the picture of a nucleus with mass M_A . Björken scaling allows to describe nuclear dynamics by the Parton Distribution Function (PDF) $F_2^A(x_A)$ which depends on the Lorentz invariant Björken variable $x_A \equiv Q^2/(2M_A\nu)[1]$. Generally the SF depends also on the resolution Q^2 which is particularly important for $x_A < 0.01/A$ where nuclear shadowing takes place. Shadowing should be included in any treatment of the EMC effect. However shadowing is described[2–4] as a multi-scattering process with diffraction between different nucleons. If the Momentum Sum Rule (MSR) has to be analyzed, the simple convolution of nucleon PDF with nuclear distribution preserves the Longitudinal Momentum (LM) of this parton system. For $x_A > 0.1/A$ we know[2, 4–6] that nuclear shadowing is unimportant.

In the light cone formulation[5, 6], x_A corresponds to the nuclear fraction of quark LM $k^+ = k^0 + k^3$ and is equal (in the nuclear rest frame) to the ratio $x_A = k^+/M_A$. But the composite nucleus is made of hadrons which are distributed with longitudinal momenta p_h^+ , where $h = N, \pi, \dots$ stands for nucleons, virtual pions, In the convolution model[5, 6] a fraction of parton LM x_A in the nucleus is given as the product $x_A = x_h * y_h$ of fractions: parton LM in hadrons $x_h \equiv Q^2/(2M_h\nu) = k^+/p_h^+$ and LM of hadrons in the nucleus $y_h = p_h^+/M_A$. The nuclear dynamics of given hadrons in the nucleus is described by the distribution function $f_h(y \equiv y_h)$ and PDF $F_2^h(x \equiv x_h)$ describes its parton structure. Let us re-

member that there are two different scales of interaction: long range nuclear scale which forms hadron distribution functions in nuclear matter and a much shorter parton scale which is responsible for their PDF's. In the convolution model restricted to nucleons and pions (lightest virtual mesons) the nuclear PDF F_2^A is described by the formula:

$$F_2^A(x_A) = \int y dy \int dx \delta(x_A - xy) (f_N(y) F_2^B(x) + f^\pi(y) F_2^\pi(x))$$

$$f_N(y) = \int \frac{d^4p}{(2\pi)^4} \delta(y - \frac{A(p^0 + p^3)}{M_A}) \text{Tr} [\gamma^+ S(p, P)] \quad (1)$$

where F_2^π and F_2^B are the parton distributions of the virtual pion and bound nucleon, P is the total four momentum of the nucleus. The function $S(p, P)$ is proportional to the connected part of the nuclear expectation value of the nucleon Green's function [7] and the trace is over the Dirac and isospin indices. Both the quark and nucleon distributions in the the formula (1) are manifestly covariant and can be expressed by Green's functions[5]. The single nucleon Green's function $S(p, P)$ in nuclear medium is given for example in [7, 8]:

$$S(p, P) = -i (\gamma \cdot (p - U_V) + M_N^*) \times \quad (2)$$

$$\left[\frac{1}{(p - U_V)^2 - M_N^{*2} + i\epsilon} + \frac{i\pi\theta(p_F - |\mathbf{p}|)}{E^*(p)} \delta(p^0 - E^*(p) - U_V^0) \right]$$

where

$$E^*(p) \equiv \sqrt{M_N^{*2} + \mathbf{p}^2}. \quad (3)$$

The values of vector $U_V = g_V(V_0, \mathbf{0})$ and scalar $(M_N^* - M_N)$ potentials are discussed in three specific mean-field models[7, 9, 10]. In the nucleus rest frame $V^- = V^+ = V^0$ for the expectation values of vector meson fields. The connected part (second term) of (3) inserted into Eq. (1)

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for $f_N(y)$ gives, after taking the trace and using the delta function to integrate over p^0 , the result which can be simplified in the RMF by using the Hugenholtz-van Hove theorem to the form[11]:

$$f(y) = \frac{4}{\rho} \int_{|p| < p_F} \frac{S_N(p) d^3 p}{(2\pi)^3} \left(1 + \frac{p_3}{E_p^*}\right) \delta(y - p^+/\varepsilon_A) \\ = \frac{3}{4} \left(\frac{\varepsilon_A}{p_F}\right)^3 \left[\left(\frac{p_F}{\varepsilon_A}\right)^2 - \left(y - \frac{E_F^A}{\varepsilon_A}\right)^2 \right], \quad (4)$$

Here the nucleon spectral function was taken in the impulse approximation: $S_N = n(p)\delta(p^0 - (E^*(p) + U_V))$, $\varepsilon_A = M_A/A$. E_F is the nucleon Fermi energy and y takes the values given by the inequality $(E_F^A - p_F)/\varepsilon_A < y < (E_F^A + p_F)/\varepsilon_A$. The Bjorken scaling in x_A corresponds in light cone dynamics to the scaling with respect to the product xy . This scaling assume that due to relativistic contraction we can in first approximation neglect the qq and NN interaction in such a deep inelastic parton structure.

The MSR for the nucleonic part is sensitive to the Fermi energy as can be seen from the integral:

$$\int dy y f_N(y) = \frac{E_F^A}{\varepsilon_A} \quad (5)$$

Thus the nucleonic part of MSR gives a factor E_F^A/ε_A which is equal to 1 at the saturation point.

$$\int F_2^A(x) dx = \int dy y f_N(y) \int F_2^B(x) dx = \int F_2^B(x) dx \quad (6)$$

To describe the "EMC effect" at the saturation where $\frac{E_F^A}{\varepsilon_A} = 1$ the nucleon PDF should be gradually smaller [8, 12, 13] in medium for Bjorken x greater than the single nucleon resolution $x_L \sim 1/M_N \rho^{\frac{1}{3}} \sim 0.4$. Consequently with eq.1 the average pion excess $\langle p_\pi^+ \rangle$ is given by the difference:

$$\frac{\langle p_\pi^+ \rangle}{M_A} = \int F_2^A(x_A) dx_A - \int F_2^N(x) dx \quad (7)$$

$$= \int_{x_L}^1 (F_2^B(x) - F_2^N(x)) dx \quad (8)$$

and is sufficiently small [8, 14, 15] ($\sim 1\%$) to describe also the nuclear Drell-Yan reactions.

However in our interacting system part of nucleons occupy states above the Fermi level. Therefore our formula (4) and MSR should be treated as the first approximation. Alternatively the mean field scenario should be supplemented by neutron-proton Short Range Correlation (SRC) which has the remarkably similar A dependence as the EMC effect[16]. On the other hand the good description of the nuclear Drell-Yan reaction with the sufficiently small [8, 14, 15] ($\sim 1\%$) pion admixture provides that our convolution approximation is up to few % accurate.

II. NON-EQUILIBRIUM CORRECTION TO NUCLEAR DISTRIBUTION.

For finite pressure very important is well known Hugenholtz van Hove relation between E_F^A , ε_A and pressure p (see for example [17]) which was proven in self-consistent mean field approach[18] to NM. The Fermi energy is defined as density derivative of the total nuclear energy $E^A = A\varepsilon_A$:

$$E_F^A = \frac{d}{dA} (E^A)_\Omega = \frac{d}{d\rho} \left(\frac{E^A}{\Omega} \right) = \varepsilon_A + \rho \frac{d\varepsilon_A}{d\rho} \quad (9)$$

where $A/\rho = \Omega$ gives the volume. The external pressure p_E is given by the thermodynamical relation:

$$p_E = - \left(\frac{\partial E^A}{\partial \Omega} \right)_A = \rho^2 \frac{d}{d\rho} (\varepsilon_A) \quad (10)$$

$$E_F^A = \varepsilon_A + E_{press}^A \quad (11)$$

Where the pressure dependent part $E_{press} = p/\rho$.

The condition (5) which measured the fraction of the nuclear LM carried by partons in nucleons gives for negative pressure p_E :

$$\int dy y f_N(y) = \frac{E_F^A}{\varepsilon_A} = 1 + p_E/\rho < 1. \quad (12)$$

The possible interpretation: the missing part of LM in (12) is carried by quarks localized inside attractive mesons(pions) responsible for this negative pressure.

However we are interested in high density region with positive pressure due to a nucleon repulsion. Above the saturation point the increasing pressure between nucleons starts to increase the nucleon Fermi energy E_F so MSR is broken by the positive factor E_F^A/ε_A . It spoils the mean field approximation and indicates that increasing pressure distorts the parton picture where the constituents

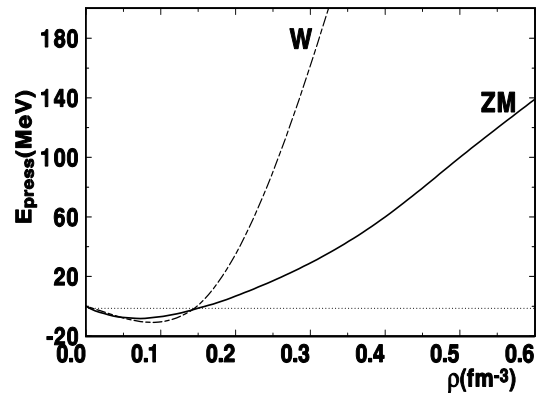


FIG. 1: E_{press} - The density dependent energy E_{press} inside NM for stiff (Walecka) and soft (ZM) EoS.

s are weekly bound in the large momentum transfer can be treated as free. Treating nucleons as bags the finite pressure will influence their surfaces [19–21]. To compensate this the nucleon PDF in the nuclear medium has to be adjusted. In particular it will be shown how the nucleon mass in medium should start to decrease in response to positive pressure. In the paper we show how such a relation change the EoS. Consider the nuclear pion contributions above the saturation point. The average distances between nucleons are smaller, therefore parameter x_L (8) increase with density. It means that the room for nuclear pions given by (8) will be reduced for higher densities. Also the pion effective cross section is strongly reduced at high nuclear densities above the threshold in $N + N = N + N + \pi$ reaction calculated in Dirac-Brueckner approach [22] (also with RPA insertions to self energy of N and Δ [23] included). Therefore for positive pressure the nuclear pions carry substantially less than 1% of the nuclear LM and dealing with a non-equilibrium correction to the nuclear distribution (1) we will restrict considerations to the nucleon part without additional virtual pions between them

$$\frac{1}{A} F_2^A(x) dx = \frac{1}{A} \int dy f_N(y) \int F_2^N(x/y). \quad (13)$$

The eventual admixture of additional pions makes the violation of longitudinal momentum even stronger. The Equation of State (EOS) for NM has to match the saturation point with compressibility $K^{-1} = 9\varrho^2 \frac{d^2 E}{d\varrho^2} \frac{1}{A}$ but then the behavior for higher densities is different for different RMF models. Generally, the choice of initial Lagrangian in the nuclear RMF models and the dependence of nucleon masses from density is not unique. We compare here two extreme examples: stiff model of Walecka [7] and soft non-linear Zimanyi-Moszkowski (ZM) model [10, 24] (see also [25]). For linear coupling in the standard Walecka model at saturation density of NM compressibility is too large ($K^{-1} \simeq 560$ MeV). The Moszkowski-Zimanyi [10] model produces the very soft EoS (see marked dotted line in fig.2) with a good value of $K-1 = 225$ MeV. The energy $E_{press} = p_E/\varrho$ shown in Fig.1 calculated for those two models reflects courses of the nuclear EOS. We see from that the E_{press} is relatively small in the ZM model (for $\varrho = 0.3 fm^{-3}$ only 25 MeV in comparison to 150 MeV in Walecka model). In both models, two coupling constants of the theory are fixed by the empirical saturation density. The additional (to the Walecka model) derivative coupling of the nucleon to the scalar field enables to re-scale the fermion wave function, and interprets the new, density dependent nucleon "Dirac" mass M_{ZM} [10, 24, 26]. It starts to decrease with density, and at the saturation point reaches 85% of free nucleon mass. But from the point of view of our description of the EMC effect [1, 12, 14], this means that 15% of nucleon longitudinal momentum is missing, and therefore has to be compensated by the enhanced pion cloud. Such a big pion cloud, however, spoils the description of nuclear Drell-Yan experiments [8, 12, 15] which measure

the sea quark enhancement. There is no evidence for a such huge enhancement in the EMC effect for $x = 0.2$. At the density $\rho = 3\rho_0 = 0.48 fm^{-3}$ $M_{ZM} \simeq 0.73 M_N$.

Our approach has different attitude, the nucleon mass begin to decrease above the saturation point. The growing pressure between nucleons starts to increase the E_F (5) and consequently the sum rule (6) is broken by the factor $E_F/\varepsilon_A > 1$. To compensate this factor which increases the longitudinal momentum (6) of nuclear partons, the the nuclear distribution and the nucleon PDF in the nuclear medium has to be changed. For good estimate, in order to proceed without new parameters, assume that the changes of PDF will be included through the changes of nucleon mass M_N in NM like in a QCD motivated bag model [20]. Multiplying the argument of the PDF by a factor E_F/ε_A the PDF will be squeezed towards smaller x and the total fraction of longitudinal momentum will be smaller by a factor ε_A/E_F^A :

$$\int_0^1 F_2^N \left(\frac{E_F^A}{\varepsilon_A} x_N \right) dx_N = \frac{\varepsilon_A}{E_F} \int_0^{\frac{E_F^A}{\varepsilon_A}} F_2^N(x) dx \cong \frac{\varepsilon_A}{E_F^A} \int_0^1 F_2^N(x) dx \quad (14)$$

Here in the integral we neglect the small contributions from $x > 1$ region originated from NN correlations.

Now, with the help of Eq.(5) and Eq.(14), the nuclear MSR is satisfied:

$$\begin{aligned} \frac{1}{A} \int F_2^A(x_A) dx_A &= \int dy y f_N(y) \int F_2^N \left(\frac{M_N}{M_{med}} x_N \right) dx_N \\ &\cong \frac{E_F^A}{\varepsilon_A} \frac{\varepsilon_A}{E_F^A} \int F_2^N(x) dx = \int F_2^N(x) dx. \end{aligned}$$

This means that quarks in the nucleus carry the same fraction of longitudinal momentum as in bare nucleons.

Such a scaling of Bjorken ratio x_N by the factor $\frac{E_F^A}{\varepsilon_A}$ in the integral (14) changes the total sum of the quark longitudinal momenta $k_{Ni}^+ = k_{Ni}^0 + k_{Ni}^3$ proportional to the quark energy in the nucleon (nucleon mass M_N). Consequently the nucleon mass in medium M_{med} for $\varrho \geq \varrho_0$ will decrease with the density by the gradually decreasing factor ε_A/E_F^A :

$$M_{med} = \frac{\varepsilon_A}{E_F^A} M_N = M_N / \left(1 + \frac{\varrho \frac{d}{d\varrho} (\varepsilon_A)}{\varepsilon_A} \right) \simeq M_N \left(1 - \frac{p_E}{\varrho \varepsilon_A} \right) \quad (15)$$

The generalized Hugenholtz van Hove theorem can be obtain [27] for different kinds of baryons in matter, for example S strange baryons (Σ, Λ) in a nuclear matter [28] with total baryon number - $(A+S)$. Analogously to Eq.11 a sum of all Fermi energies, including Fermi energies of strange baryons E_F^S , is equal to the total energy $E^{A+S} = M_{A+S}$ with analogous pressure corrections:

$$\begin{aligned} AE_F^A + SE_F^S &= M_{A+S} + (A+S) p_E / \varrho \\ \text{with } p_E &= - \left(\frac{\partial E^{A+S}}{\partial \Omega} \right)_{A+S} \end{aligned} \quad (16)$$

The corrections (15) depend only from pressure and total energy of the system therefore the similar medium correction based on Eq.16 will apply to masses of different barions, like strange Λ and Σ , in NM.

Let us discuss these results in the simple bag model[29] where the nucleon in the lowest state of $n = 3$ quarks is a sphere and its energy E_{Bag} is given as a function of radius R with phenomenological constants - ω_0, Z_0 and B :

$$E_{Bag} = \frac{n\omega_0 - Z_0}{R} + \frac{4\pi}{3}BR^3 \quad (17)$$

This formula was obtained by studying the free quark flow described by the stress tensor $T_D^{\mu\nu}$ is given on bag surface along the vertical n^μ to this surface by:

$$n_\mu T_D^{\mu\nu} = \frac{1}{2} \frac{\partial}{\partial x_\nu} \sum \bar{q}_a(x) q_a(x) \quad (18)$$

where $q_a(x)$ denotes the quark field with given flavor and color. Because the $\bar{q}_a(x) q_a(x) = 0$ on the surface its derivative must lie along the normal and define the intrinsic pressure p_D .

$$n_\mu T_D^{\mu\nu} = n_\nu p_D \quad (19)$$

The pressure condition for p_D can be obtained from a usual criterium (11):

$$(\partial E_{Bag} / \partial \Omega_{Bag})_n = 0 \quad (20)$$

However in a compressed medium the pressure p_D generated by free quarks inside the bag is balanced at the bag surface[29] not only by a intrinsic confining force represented by the bag constant B but additionally by external pressure p_E generated by elastic collisions with other hadrons[19, 20, 30]. In this way the residual strong interaction with all other quarks localized outside the given bag is taken into account:

$$p_D = B + p_E \quad (21)$$

Finally using the equation Eq.(20) and resulting spherical solution for the bag radius we obtain:

$$\frac{3\omega_0 - Z_0}{R^2} = 4\pi(B + p_E)R^2$$

$$R = \left[\frac{3\omega_0 - Z_0}{4\pi(B + p_E)} \right]^{1/4} \quad (22)$$

At the saturation $p_E = 0$ and the bag constant B can be determined by the value of the nucleon radius $R \simeq 0.6 fm$. Above the saturation point the R increases[20, 21]. The main reason: the external pressure was not taken into account and the value of B wrongly determines the increasing bag radius R (nucleon swelling). Here we include explicite the external pressure which increases when the bag constant decreases. When the sum $(B + p_E)$ weakly depends from density,

then the bag radius will remain, according to (22) approximately constant with the increasing nuclear density ϱ . In fact such a assumption is confirm in the calculation of [20] where an decreasing of the B constant from a saturation density ρ up to 3ρ is accompanied by similar increase of the external pressure p_E . The QMC model in medium[21, 31] allows to find how the nucleon radius is changing with density and for the derivative coupling model with ZM[10] Lagrangian the nucleon radius remains almost constant[31] up to density $\varrho = 10\varrho_0$.

If for $\varrho > \varrho_0$ $B(\varrho) + p_E(\varrho) \simeq B(\varrho = \varrho_0)$ that the nucleon energy in the rest frame can be written as:

$$E_{Bag} = 4\pi R^3 \left[\frac{4}{3}(B + p_E) - \frac{p_E}{3} \right] \simeq E_{Bag}^{p_E=0} \left[1 - \frac{p_E}{\varrho E_{Bag}^{p_E=0}} \right] \quad (23)$$

with $\varrho \simeq 1/\Omega_{Bag} > \varrho_0$ where the volume of a space between nucleons above the saturation density ϱ_0 was neglected. E_{Bag} differs from the the nucleon mass by the c.m. correction [21] to the nucleon parton model. Now we can compare the expressions (15) with (23) for the nucleon mass in medium. The pressure corrections essentially given by $E_{press} = (p/\varrho)$ are similar, especially when the nucleon radius and $(B + p_E)$ remains constant[31] in the bag (23) model. Such a solution of $R(\varrho)$ which is constant with ϱ testifies the properly soft EoS; i.e. ZM model with $K^{-1} \simeq 250$ MeV showed in Figs.1,2. The pressure corrections to the masses of strange baryons will be different because the important ingredient the mass of the strange quark (300MeV) remains constant and the reduction (23) concerns the remaining massless part. In effect the mass difference between strange baryons and nucleons will increase slightly. Also the best candidate[28, 32] to strange admixture the Σ^- has the repulsive[33] $\Sigma - NM$ interaction which also shifts the possible strangeness admixture to higher density[32]. Therefore we present now the calculation restricted to nucleons as a good estimate for densities from ϱ_0 to $3\varrho_0$. The full calculation which will include admixture of strangeness will be published elsewhere.

Now our explicit mass dependence from density, energy ε_A and energy derivative (15) is plugged into the following standard Walecka RMF equations[7] for energy per nucleon ε_A and effective mass M^* :

$$\varepsilon_A = C_1^2 \rho + \frac{C_2^2}{\rho} (M_{med} - M_{med}^*)^2 + \frac{\gamma}{(2\pi)^3 \rho} \int d^3 p \sqrt{(p^2 + M^{*2})}$$

$$M_{med}^* = M_{med} - \frac{\gamma}{2C_2^2 (2\pi)^3} \int d^3 p \frac{M_{med}^*}{\sqrt{(p^2 + M^{*2})}} \quad (24)$$

where γ denotes level degeneracy ($\gamma = 2$ for neutron matter) and two (coupling) constants: vector C_v^2 and scalar C_s^2 , were fitted[7, 9] at the saturation point of nuclear matter (in the formula $2C_1^2 = C_v^2/M_N^2$ and $2C_2^2 = M_N^2/C_s^2$).

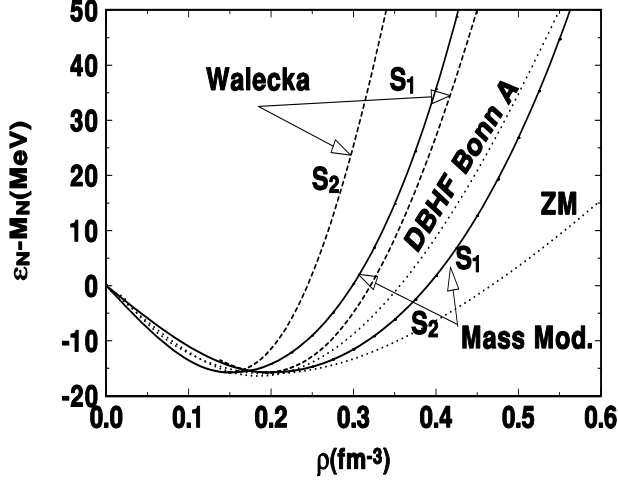


FIG. 2: The nucleon energy $\varepsilon_A - M_N$ as a function of NM density for two RMF models; scalar-vector Walecka (dot lines) and our Modified Mass approach (solid). Both RMF models are calculated for two parameterizations: S_1 version[7] ($\rho_0 = .19 fm^{-3}$) and version[9] S_2 ($\rho_0 = .16 fm^{-3}$). Results for full DBHF[36] (dotted marked line) calculation using Bonn A NN interaction are displayed for comparison, also nucleon energy in ZM model[10] is in the plot (dotted marked line).

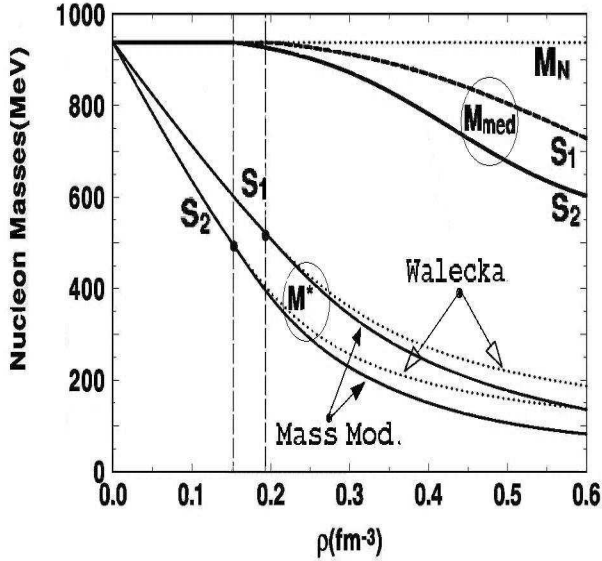


FIG. 3: Constant nucleon mass M_N (Walecka) and density dependent mass M_{med} from our "Mass Mod." model. Also effective mass M^* in these approaches. Both models are calculated for S_1 and S_2 parametrization.

In the Walecka model $M_{med} = M_N$. In our model the finite pressure corrections to M_{med} (15) convert the recursive equation (24) to a differential equation above the saturation density ρ_0 in general form:

$$f(\varepsilon_A, \frac{d}{d\rho}(\varepsilon_A)) = 0 \text{ for } \rho \geq \rho_0 \quad (25)$$

Note that equation (24) is obtained from the energy-momentum tensor for the model Hamiltonian with constant nucleon mass[7]. Here we assume that the same equation with medium mass M_{med} will be satisfied. It should be a good approximation, at least not very far from the saturation density. The pressure p is obtained from the thermodynamic relation (11).

The final results were obtained by solving¹ numerically differential recursive equations (25), starting from standard solutions of Eq.(24) at the saturation density for two versions of the Walecka model: first version S_1 [7] have a minimum at $\rho_0 = .19 fm^{-3}$ (parameters $C_v^2 = 195.9, C_s^2 = 267.1$) and second version S_2 [9] ($\rho_0 = .16 fm^{-3}, C_v^2 = 273.8, C_s^2 = 357.4$). They are displayed in Figs.(2-4). In Fig.2 our values of energy per nucleon ε_A calculated for two versions S_1 and S_2 are denoted by solid lines (Mass Mod.) and solutions of the ordinary Walecka model with constant mass, denoted by dashed lines, are presented for comparison. Our EoS's are generally much softer - from the unrealistic value of $K^{-1} = 560 MeV$ for the Walecka model (S_1) to the reasonable $K^{-1} = 290 MeV$ obtained in our model. Below saturation density the solutions are of course identical (solid lines). The nucleon mass M_{med} and the nucleon effective mass M^* (both in Walecka and in our (Mass Mod.) model) are compare in Fig.3. The decrease of M_{med} is directly connected with the change of nucleon SF for positive pressure in NM. The decrease of M_{med} above saturation is directly connected with the change of nucleon PDF for positive pressure in NM. In Fig.4 we plot the pressure against NM density. Our results are compared here with a semi-experimental estimate[35] from heavy ion collisions and indeed correct (solid lines) Walecka results (dashed) quite well making the EoS significantly softer. For S_1 parametrization our model in Fig.4 fits well the allowed region in neutron matter. For nuclear matter we have good course of EOS below density $\rho = 5 fm^{-3}$. In fact, for this density, the (partial) de-confinement is expect which will change EOS above the phase transition[37]. Our results are close (slightly below) DBHF results (dotted line). Therefore it is interesting how the realistic NN calculations with off-shell effects correspond with the modification of the nucleon PDF in medium induced by MSR on the parton level.

¹ It is important to mention that in solutions of Eqs.(24,25) the Fermi energy from definition Eq.(9) has a different value from the usual form $E_F = \sqrt{M^{*2} + p_F^2} + U_v$ used in Eq.(4). The discrepancy vanish near the saturation density, increases with the density and reach the 15% of the total vector repulsion in Eq.(24). Similar problems are well known in the saturated NM[34] and are connected with the proper choice of single particle potential which in our case should be adjusted to the changes of nucleon mass. This discrepancy can be removed, here e.g. by the less repulsive momentum dependent vector potential for nucleons however such a correction has no influence on presented mass changes of the nucleon and our final results.

For example, in DBHF method, there are additional corrections[36] from self energy which diminish the nucleon mass with density.

III. CONCLUSION

We have presented a simple, parameter free, hybrid model of nuclear matter, which examines the influence of the parton structure on nuclear EoS. It was shown that there is connection between stiff EOS and strong viola-

tion of longitudinal momentum conservation carried by quarks. In our approach we propose a scaling of the nucleon mass which modify the nucleon PDF in order to satisfy the longitudinal momentum sum rule for the partons in NM above the saturation density. Our pressure correction reduce the violation of MSR from 50% in the linear Walecka model to 10% in our model. This modify the EoS for nuclear and neutron matter making it softer, close to semi-empirical analysis[35] and close to DBHF calculation with a realistic Bonn A potential[36]. The similar mass corrections concern all baryons in medium and will be present e.g. in the strange nuclear matter. Other features of the Walecka model, including a good value of the spin-orbit force remain in our model unchanged. Our results suggest pressure corrections above the saturation density to any RMF model with constant nucleon mass and unmodified parton PDF (without a properly normalized nucleon longitudinal distribution). Finally we argue that the nucleon radius or the confining region should remains constant with density otherwise the pressure correction to the nucleon mass will be different in Bag model Eq.(23) versus parton model Eq.(15). The alternative: increasing confining region in the Bag model which corresponds eventually to the increasing number of intermediate pions(mesons) for a high density will give a very stiff EoS; rule out when studying star properties (mass-radius constraint), especially the mass of PSR J16142230 the most massive known neutron star[38] recently discussed[37] in the application to compact and neutron stars. Partial support of the Ministry of Science and Higher Education under the Research Project No. N N202046237 for is acknowledged.

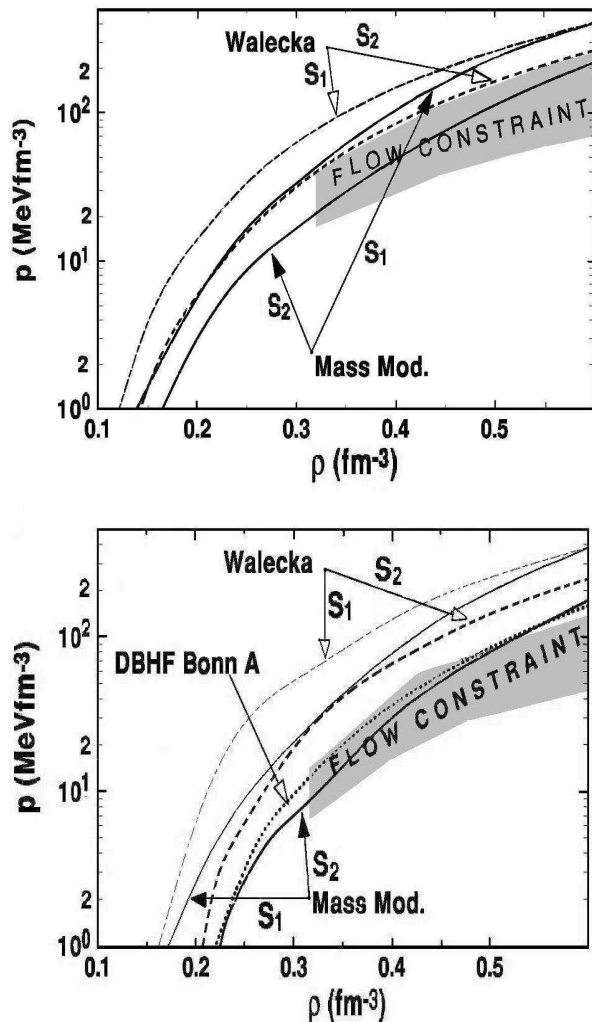


FIG. 4: Pressure for neutron matter (upper plot) and for nuclear matter (bottom) as a function of density for two most frequent pure scalar-vector parameterizations of Walecka RMF model: version[7] S_1 ($\rho_0 = .16 fm^{-3}$) and version[9] S_2 ($\rho_0 = .16 fm^{-3}$) are denoted by dashed lines. Our (Mass Mod.) modification are denoted by solid lines. The area denoted by "flow constrain" taken from[35] determined the allowed course of EoS, using analysis which extracted from matter flow in heavy ion collisions the high pressure obtained there. The DBHF (ref.[36]) calculation with Bonn A interaction are shown as dotted line for comparison.

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